

RENORMALIZATION-GROUP FLOW ANALYSIS OF MESON CONDENSATIONS IN DENSE MATTER

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ABSTRACT

We present a renormalization-group (RG) flow argument for s-wave kaon condensation in dense nuclear-star matter predicted in chiral perturbation theory. It is shown that it is the *relevant* mass term together with *any* attractive interaction for the kaon in medium that triggers the instability. We show that a saddle point of multi-dimensional RG flow can imply a phase transition. Pion condensation is also analyzed along the same line of reasoning.

The idea of renormalization group (RG) has been used extensively both in condensed matter physics and particle physics, especially for the critical phenomena and various situations involving scaling behavior. Recently, Shankar[1] and Polchinski[2] showed that one can use the RG idea even for the phenomena including a scale, such as the mass gap, as in BCS superconductivity and charge density wave etc. The key point is that these phenomena could be understood as an instability from the Fermi liquid identified as the fixed point of the RG flow. They showed for the BCS case as an example that when two incoming momenta sum to zero, the corresponding four-Fermi interaction is *marginally relevant* ^{#1} in the RG sense, so that the interaction causes an instability that pushes away the system from the Fermi liquid. The new insight gained in this approach is that one can identify in a clear and simple way the dynamics and kinematics that lead to the phase transitions. In this paper, we extend this approach to the condensation of negatively charged kaons (K^-) in dense nuclear medium as in neutron stars by considering the role of the quadratic term in the effective potential entering in kaon-nucleon interactions. (The K^+ meson does not condense and the neutral kaons K^0 and \overline{K}^0 are not relevant in neutron stars.)

Recently Lee *et al* [3, 4] have shown by chiral perturbation theory (χPT) treated to *in-medium* two-loop order (corresponding to next-to-next-to leading order) that kaons can condense in dense nuclear-star matter at a matter density $\rho \lesssim 4\rho_0$ where ρ_0 is the normal nuclear matter density. However to the order considered, many terms are involved and it is not transparent which mechanism is in action for triggering the condensation process.

In this note, we present a renormalization group flow analysis to show what drives the process of kaon condensation and in particular to indicate the basic mechanism involved. The conclusion is that kaons must condense in s-wave, although the analysis cannot give the critical density.

For the purpose of elucidating the basic concept, we find it sufficient to study a toy model which we believe captures the essence of the physics involved in kaon condensation. The corresponding action, S , can be decomposed into three parts: S_K for the free kaon, S_N for the nucleon and S_{KN} for kaon-nucleon interactions. We assume that nucleons in nuclear matter are in Fermi-liquid state with the Fermi energy μ_F and the Fermi momentum k_F . This state might arise from chiral Lagrangians as some sort of “Q-balls” or nontopological solitons. (For a discussion on this, see ref.[5].) For our purpose, it is crucial that the nuclear matter arises as a Fermi liquid[6]. Defining ψ as the nucleon field fluctuating around the Fermi surface such that the momentum integral has a cut-off Λ_N ,

$$k_F - \Lambda_N < |\vec{k}| < k_F + \Lambda_N, \quad (1)$$

^{#1}Throughout this article, we put in *italic* the terms “relevant,” “marginal,” and “irrelevant” when used in the RG sense.

the action in the nucleon sector can be written, schematically, in the form

$$S_N = \int d\epsilon d^3k \psi^\dagger (\epsilon - \epsilon(\vec{k})) \psi + g \int (d\epsilon d^3k)^4 \psi^\dagger \psi^\dagger \psi \psi \delta^4(\epsilon, \vec{k}) \quad (2)$$

where $\epsilon(\vec{k})$ is the nucleon energy measured as usual relative to the chemical potential μ_F and the δ function assures the overall energy-momentum conservation. The second term with a coupling constant g is a generic residual Landau-Migdal four-Fermi (nucleon) interaction^{#2} which is *marginal*, so the normal nuclear matter is a fixed point, which will be called “Fermi-liquid fixed point” as in condensed matter physics. This is an immediate application of the recent idea developed by Shankar[1] and Polchinski[2, 7] to nuclear matter.

We focus on kaons that are in chemical equilibrium with electron gas in high density. So the kaon chemical potential – denoted μ_K – is the same as that of the electron. This follows from charge conservation. In this situation, it is more convenient to measure the kaon frequency relative to the kaon chemical potential. Furthermore the kaon fluctuating momentum, \vec{q} , is much smaller (*i.e.*, in s-wave) than the kaon chemical potential, so the kaon field K can be taken to be essentially non-relativistic. (This is most likely to be valid for kaon condensation in compact-star matter where the kaon chemical potential is substantial but may not be valid in dense symmetric nuclear matter appropriate for heavy-ion collisions.)

Therefore what is relevant is the small kaon energy fluctuation around the chemical potential. These properties of the kaon can be simply incorporated by redefining the kaon field as $K \rightarrow \Phi e^{-i\mu_K t} / \sqrt{2\mu_K}$. The action for the kaon is

$$S_K = \int d\omega d^3q \Phi^*(\omega, \vec{q}) \left(\omega - q^2/2\mu_K \right) \Phi(\omega, \vec{q}) - \int d\omega d^3q \tilde{M} \Phi^* \Phi \quad (3)$$

which follows from a relativistic Lagrangian by keeping terms of order $1/\mu_K$. The quadratic coupling (“mass”), \tilde{M} , in the second term is $\frac{M_K^2 - \mu_K^2}{2\mu_K}$. As with all effective theories, we need a cut-off for the fluctuation. The momentum cut-off for the kaon is defined by Λ_K , $|\vec{q}| < \Lambda_K$, which cannot be larger than μ_K . For simplicity, we shall not distinguish between Λ_K and Λ_N and denote them generically Λ .

For kaon-nucleon interactions, we focus on s-wave kaons. (P-wave kaons do not figure directly in condensation phenomena for the same reason as the absence of p-wave pion condensation as mentioned below.) The s-wave kaon-nucleon interaction involved can be summarized in a simple form

$$S_{KN} = \int (d\omega d^3q)^2 (d\epsilon d^3k)^2 h \Phi^* \Phi \psi^\dagger \psi \delta^4(\omega, \epsilon, \vec{q}, \vec{k}). \quad (4)$$

The important quantity here is the coefficient h . In chiral Lagrangians, this coefficient subsumes several terms of varying importance, including energy dependence. For instance,

^{#2}Because of the flavor degree of freedom, there are in general four independent four-Fermi interactions for nuclear interactions. They are all *marginal*.

terms of higher chiral order but quadratic in the nucleon field also take this form. Specifically in chiral perturbation theory[3, 4]^{#3}, the h includes^{#4} the leading chiral order term $[O(Q)]$ of the form $\sim \frac{\omega}{f_\pi^2} K^\dagger K N^\dagger N$ and the next-to-leading-order terms $[O(Q^2)]$, the so-called sigma term, $\sim \frac{\Sigma}{f_\pi^2} K^\dagger K \bar{N} N$ and a higher derivative term, $\sim \omega^2 K^\dagger K \bar{N} N$. (Here K is the doublet kaon field, N the doublet nucleon field and ω is the kaon frequency that equals μ_K at condensation.) In the case of s-wave kaon condensation, what we are interested in is the flow of the quadratic term (*i.e.*, the “mass” term) in the effective potential under the RG transformation.

Now we have a well defined system of strongly interacting kaons and nucleons as a system with a possible ground state configuration determined by the chemical potentials of the kaon and the nucleon. This configuration is supposed to result from many-body effects of the nuclear matter. The interactions between fluctuating fields around the ground-state configurations define the toy model, the action of which is taken to be of the form

$$\begin{aligned}
S &= S_K + S_{NK} + S_N \\
&= \int d\omega d^3q \Phi^*(\omega, \vec{q}) \left(\omega - q^2/2\mu_K \right) \Phi(\omega, \vec{q}) - \int d\omega d^3q \tilde{M} \Phi^* \Phi \\
&\quad + \int (d\omega d^3q)^2 (d\epsilon d^3k)^2 h \Phi^* \Phi \psi^\dagger \psi \delta^4(\omega, \epsilon, \vec{q}, \vec{k}) \\
&\quad + \int d\epsilon d^3k \psi^\dagger (\epsilon - \epsilon(k)) \psi + g \int (d\epsilon d^3k)^4 \psi^\dagger \psi^\dagger \psi \psi \delta^4(\epsilon, \vec{k}). \tag{5}
\end{aligned}$$

We now perform the RG analysis following Shankar[1] and Polchinski[2, 7]. The low-energy effective theory with a cut-off $s\Lambda$ with $s < 1$ can be obtained by integrating out the high frequency modes. The stability of the system can then be determined by computing the RG flow of the coupling constants of the interaction terms following this mode-elimination procedure.

The scaling law of the kaon field is *defined* by requiring the kinetic term of eq.(5) to be invariant under scaling $\omega \rightarrow s\omega$ and $q \rightarrow \sqrt{s}q$, which is imposed to make ω and $q^2/2\mu_K$ scale in the same way. Since the μ_K is fixed by the interaction with the electron, we can keep it scale-free. The scaling dimension of Φ then comes out to be^{#5} $[\Phi] = -7/4$ which follows simply from the scaling of the integration measure $[d\omega d^3q] = 5/2$ and the invariance condition $[d\omega d^3q] + [\omega] + 2 \times [\Phi] = 0$. The second term (“mass” term) is not invariant but

^{#3} Our strategy in applying effective chiral Lagrangians to our many-body problem is as follows: We assume that we first retain suitable leading chiral order terms corresponding to a mode elimination in the “elementary-particle” space and then apply the mode elimination around the Fermi sea *afterwards*.

^{#4}The s-wave $K^-p\Lambda(1405)$ coupling that plays a crucial role in the threshold properties of kaon-nucleon interactions and also in K-mesic atoms is not included in the model we are considering. It turns out to play a minor role (as an *irrelevant term*) in the RG-flow argument, lending a support to the similar finding in chiral perturbation theory of [3, 4].

^{#5} $[\mathcal{O}]$ stands for the scaling dimension of the object \mathcal{O} .

relevant, $[d\omega d^3q] + 2 \times [\Phi] = -1$, which means that the “mass” term grows as the fast modes are integrated out. But this is not the entire story since there is a contribution from the interaction term when fast nucleon modes in $s\Lambda < k < \Lambda$ are eliminated. To see what that does, we have to determine the scaling behavior of the nucleon system. We take the well-known procedure of the “scaling toward the Fermi surface” where only the component perpendicular to the Fermi surface is scaled. We have the scaling $[d\epsilon d^3\vec{k}] = 2$, $[\delta^4(\epsilon, \vec{k})] = -2$. Again keeping the kinetic part – the fourth term of eq.(5) – invariant, we get the scaling rule for the nucleon, $[\psi] = -3/2$, $[d\epsilon d^3\vec{k}\psi] = 1/2$. This allows us to calculate the scaling of the third term of eq.(5)

$$2 \times ([d\omega d^3\vec{q}\Phi] + [d\epsilon d^3\vec{k}\psi]) + [\delta(\omega, \epsilon)] + [\delta(\vec{q}, \vec{k})] = 1/2.$$

This shows that the kaon-nucleon four-point interaction is generally *irrelevant*. The question we wish to address now is whether the scaling law can be modified by radiative (loop) corrections of the sort discussed recently by Polchinski[7] and Nayak and Wilczek[9]. A simple analysis shows that up to one loop, there are no corrections that change the tree-level scaling behavior. For example, the correction from Fig. 1b is *irrelevant* since the overlapping kinematical region for the kaon and nucleon shrinks as s decreases. This completes the counting of the naive scaling dimensions of the terms involved. At this point one might conclude that the common wisdom for s-wave kaon condensation cannot be supported by the RG-flow argument, since the attractive interactions involving Σ term which are supposed to drive the kaon condensation are found to be *irrelevant*. However we will show that this *irrelevant* but attractive term is what triggers the kaon condensation.

Under the RG transformation with $\Lambda \rightarrow s\Lambda$, the perturbation on the “mass” can be calculated as

$$\delta\tilde{M} = s^{-1} \left(\tilde{M} - (1-s) \frac{\Lambda h}{6\mu\pi^2} (3k_F^2 + \Lambda^2(1+s+s^2)) \right) - \tilde{M}, \quad (6)$$

where the factor s^{-1} comes from the scaling property of Φ . The one-loop RG equations can be obtained by putting $s = 1 - \delta t$ in eq. (6)

$$\frac{d\tilde{M}}{dt} = \tilde{M} - Dh, \quad (7)$$

$$\frac{dh}{dt} = -ah - Ah^2, \quad (8)$$

where, $D = \frac{3(1+\alpha^2)\alpha}{2\mu_K}\rho$, $t = -\ln s$, $\alpha = \Lambda/k_F$. The constant a can be obtained from the scaling properties of the interaction terms: For the interaction in eq.(4), $a = 1/2$. The second term on the right-hand side of eq. (7) is given by the diagram of Fig. 1a. One can

show that A in eq. (8) which can be calculated from the diagrams in Fig.1b vanishes ^{#6}. The factor $\frac{\Lambda}{k_F}$ in the coefficient D is due to the restriction imposed on the momentum region, *i.e.*, the momentum cut-off. We are considering fluctuations near the Fermi surface. The solution for the \tilde{M} -flow is given by

$$\tilde{M}(t) = (\tilde{M}_0 - \frac{Dh_0}{1+a})e^t + \frac{Dh_0}{1+a}e^{-at} \quad (9)$$

with

$$h(t) = h_0 e^{-at}, \quad h_0 \geq 0. \quad (10)$$

In this formula, \tilde{M}_0 must be positive as the boson chemical potential cannot exceed its mass.

In order to understand the physics involved, we describe the flow for general a and D in the (\tilde{M}, h) plane. For the moment, we shall forget that these quantities are fixed in our model and allow arbitrary values to a and D . In general, D is positive if the interaction is attractive, negative if repulsive and zero if there is no interaction. Also $a > 0$ if the interaction is *irrelevant*, < 0 if *relevant* and $= 0$ if *marginal*. We describe the flow of each case in Figure 2. If there were no attractive interaction h giving rise to a density-dependent term in the mass flow [eq.(7)], then the flow would have gone straight up in the mass direction in the (\tilde{M}, h) plane if we start from any initial value $\tilde{M}_0 > 0, h_0 > 0$. See Fig. 2 a,b,d,f,h. In this case, if the mass is positive at some scale, then it will remain positive at all scale, so there would be no region from which the mass could flow to negative direction. However, with inclusion of an *attractive* interaction, the mass flow becomes qualitatively different. See Fig. 2 c,e,g. Our case defined by the toy model (5) corresponds to Fig. 2c. This analysis shows us the nontrivial aspects of *irrelevant* interactions in determining the direction of “mass” flow. We also note that the Gaussian fixed point is a saddle point of the RG flow and unstable.

The interesting feature of this flow in connection to kaon condensation is that there are two types of flows depending upon the sign of the coefficient of e^t in eq.(9). For

$$D < \frac{\tilde{M}_0}{h_0}(1+a), \quad (11)$$

it flows toward $+\infty$ whereas for $D > \frac{\tilde{M}_0}{h_0}(1+a)$, it flows downward to $-\infty$. The flow eventually crosses the zero “mass” axis with the speed of the RG flow determined by the density ρ_N . Thus in the low-energy limit, the sign of the “mass” term becomes negative if the the initial point belongs to the shaded region in Fig. 2c. This signals a meson condensation as we explain below.

^{#6}Here we assume that the BCS-type instability for four nucleon coupling is suppressed. In fact there is no BCS-type instability involving strangeness flavor.

In the mean-field approach, when the effective mass of the kaon decreases and approaches μ_K , kaon condensates will develop. This may be expected from the ideal bose gas distribution

$$n(k) = \frac{1}{e^{\beta[\epsilon(k)-\mu]} - 1}, \quad (12)$$

where the ground-state occupation number diverges when the chemical potential equals the ground-state energy. What the RG analysis tells us is that this tendency for the kaon condensation will be eliminated in the low energy limit if the “mass” flows toward a point $(\infty, 0)$ in the (\tilde{M}, h) plane. Therefore the flow toward a point $(-\infty, 0)$ or crossing the zero “mass” axis can be identified as a signal for kaon condensation in the RG analysis. A similar observation can be made from the mechanism of spontaneous symmetry breaking. As mass becomes negative, the vacuum becomes unstable, and something must happen for the system to restore the stability. What happens is that the system develops a positive Φ^4 term for stability and the vacuum gets shifted to a new one with non-zero VEV, $\langle \Phi \rangle \neq 0$. This is the familiar mechanism for spontaneous symmetry breaking induced by the sign change in the mass. Therefore the “mass” flow crossing zero can be naturally associated to an instability conducive to a certain phase transition. *The key point of this paper is that the downward flow of Fig.2c can be associated with the instability leading to kaon condensation.*

An interesting point to note in the above analysis is that the main driving term for m in the RG equation is provided by the *relevant* term, *i.e.*, the mass term, but the direction of the flow is determined by the *irrelevant* interaction term. An illustrative application of this observation is that in the chiral limit, where there is no mass term, condensation cannot be driven by the remaining *irrelevant* interactions. This is consistent with the fact that, in the chiral limit, chiral symmetry prevents the flow of the mass and there cannot be any phase transition.

To understand better the role played by chiral symmetry breaking in kaon condensation, we examine how chirally symmetric terms and symmetry-breaking terms differ in their contribution to the mass or energy. To do this, we take one term each from three classes of kaon-nucleon interactions: (i) a mass term, (ii) an explicit chiral-symmetry-breaking interaction term (*i.e.*, the Σ term), $\frac{\Sigma}{f_\pi^2} K^\dagger K \bar{N} N$, and (iii) a chiral-symmetry-preserving interaction term, $c \partial_\mu K^\dagger \partial^\mu K \bar{N} N$. Notice that in our toy model above, (i) and (ii) are included in the h term. With one-loop corrections to the self-energy, the relativistic inverse propagator becomes

$$G^{-1}(k, \omega) = \omega^2 - k^2 - m^2 - \Pi(\omega, k) \quad (13)$$

where m is the meson mass, $\Pi = C(\omega^2 - k^2) - H$ with C a constant proportional to c and $H = 3 \frac{\Sigma}{f_\pi^2} (1 + (\Lambda/k_F)^2) (\Lambda/k_F) \rho_N$ is the contribution from the Σ term. Equation (13) can

be rewritten as

$$G^{-1}(k, \omega) = (1 - C)(\omega^2 - k^2 - \tilde{m}^2), \quad (14)$$

where $1 - C$ is the wave function renormalization factor and $\tilde{m}^2 = (1 - C)^{-1}(m^2 - H)$. The mass shift is

$$\delta m^2 = (1 - C)^{-1}(m^2 - H) - m^2 = (1 - C)^{-1}(Cm^2 - H). \quad (15)$$

Let us first turn off the Σ term. Then the mass correction is simply proportional to the original mass, $\delta m^2 = \frac{C}{1-C}m^2$. Thus if the original mass is small, then the chirally symmetric interaction gives a small mass correction. This is just the precise statement of PCAC. If there were no mass term to start with, of course, there would be no mass correction. In this case there would be no Σ term either. The Goldstone boson mass is protected by chiral symmetry. However if the symmetry breaking is not small, chirally symmetric interactions can contribute importantly to mass or energy. The above analysis shows that repulsive interaction ($C > 0$) increases the meson mass while attraction decreases it. When the Σ term is present, the Σ term attraction and the C term repulsion can compensate each other. For pion-nucleon interactions, the two terms effectively cancel, leaving the pion mass more or less unchanged, while in the case of kaons, the Σ term attraction wins out.

An exactly parallel RG-flow analysis could be made for s-wave *pion* condensation. As in the case of the kaon, it is the mass term that can cause instability with the $\Sigma_{\pi N}$ term determining the direction of the “mass” flow. We therefore expect a similar mass flow as in Fig. 2c, leading to s-wave condensation. However the pion is *almost* massless on the strong-interaction scale and the chiral symmetry *almost* protects its mass. Any small repulsion would counterbalance the small attraction associated with the $\Sigma_{\pi N}$ term. Thus as we know from on-shell charge-symmetric πN and π -nuclear amplitudes, the repulsive term $\sim \omega^2$ – which includes, in nuclear medium, the well-known Pauli exclusion principle effect[12] – overpowers the $\Sigma_{\pi N}$ attraction, preventing the pion mass from going to zero. More specifically, in the RG analysis, $\tilde{M}_0 - Dh \geq 0$ for the s-wave pion due to the repulsive term $\sim \omega^2$ and there is no instability toward s-wave pion condensation. This is consistent with PCAC which says that the soft-pion point – which is physical in the chiral limit – and the would-be condensation point are smoothly interpolated and hence the absence of s-wave pion condensation in the chiral limit implies its absence at any off-shell point. This is also consistent with the observations made in lattice gauge calculations (in high temperature) and QCD sum rule calculations which indicate that the pion mass does not change appreciably as temperature or density is increased. Perhaps the most significant point to note is that in two extreme limits, the chiral limit and the heavy-quark limit, bose condensations cannot take place and that kaons can condense at low enough density because the kaon is neither very light nor very heavy.

The p-wave pion condensation studied extensively since many years is different from the s-wave case from the point of view of the RG flow. As shown, in the case of the s-wave condensation, it is the fluctuation in the “mass” direction around zero meson three-momentum that affects crucially the low-energy dynamics (the role of the s-wave pion interaction is minor in the sense that the interaction itself is not responsible for the instability). The condensate so developed is spatially uniform, with the order parameter being space-independent. However the p-wave condensation involves non-zero spatial momentum, with the condensate varying in space. For this, Yukawa interactions with the meson of non-zero momentum are required. For an instability to set in, a *relevant* interaction in the appropriate channel would be necessary. However, we can show, following Polchinski [7], that the πNN Yukawa interaction, when radiatively corrected, becomes *irrelevant* and the four-Fermi interaction in the pion channel is at most *marginal*. Therefore p-wave pions do not condense. It is not clear how this reasoning is related to the mean-field argument [10] where the strength of the Landau-Migdal interaction g'_0 plays a crucial role in pushing up the critical density.

One might ask what the effect of the kaon condensation is on the fermion (nuclear) system. The fermion system involved here is a bit intricate because of the Fermi surface. Since the density that gives the Fermi momentum is a fixed and physically controlled quantity, we must compensate the change by introducing a counter term[1]. This means that the flow of the Fermi surface should be prevented by *fiat*. This is the basic difference between the kaon and nucleon mass-like terms. While the flow of the kaon mass is allowed and leads to instability, we cannot allow the flow of the Fermi surface since the density must be fixed.

Now once the kaon condensation takes place, one can decompose the kaon field into a condensed part and a fluctuating part. Focusing on the s-wave condensation only, we can write

$$K(x, t) = K_0 e^{-i\mu t} + \int \frac{d^4 k}{(2\pi)^4} \Phi(k) e^{ik \cdot x} \quad (16)$$

where $k = (\omega, \mathbf{k})$ and K_0 represents the condensed part. With eq.(16), the interaction term in the model Lagrangian, eq.(5), contains an additional term

$$\int d\epsilon d^3 k h |K_0|^2 \psi^\dagger \psi. \quad (17)$$

This means that the net effect of the kaon condensation on fermion dynamics is to change the fermion chemical potential by $\mu_F \rightarrow \mu_F + h|K_0|^2$. If we fix the density of the nucleons, then the Fermi surface is fixed, $k_F \sim n^{1/3}$. The effective fermion mass must therefore decrease since $\mu_F = k_F^2/2m_F^*$ must increase.

In this note we have discussed how the RG approach can be used for strongly interacting hadronic matter, *e.g.*, for the case where kaon condensation might take place. The flow of the kaon “mass” (*i.e.*, the quadratic coupling) is argued to be a signature for the

instability of the kaon-nuclear system. We interpret the RG flow across the zero mass axis as the instability that leads to the s-wave kaon condensation. An analogy to the familiar spontaneous symmetry breaking (SSB) mechanism has been used, where the negative mass term flows due to the mode elimination from the Φ^4 term and eventually crosses zero. At this point the vacuum changes into the “true” vacuum with $\langle \Phi \rangle \neq 0$. We should emphasize that the mechanism we are suggesting is so general that it is independent of the detailed structure and of the number of interaction terms. What we are showing here is that in *any* interacting boson-fermion systems, s-wave bose condensation must take place if there is an attractive interaction between fermion and *massive* boson and if the density is high enough. Similar RG flow analyses can be made for cases where there are more than two coupling constants. Although our discussion is based on a simplified model, it contains, however, enough generic feature that encapsulates the basic physics implied by the underlying symmetry of the system, namely, chiral symmetry and its breaking[5].

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FIGURE CAPTIONS

Figure 1: One-loop corrections.

- a) Mass correction (to \tilde{M}) due to $K^\dagger K N^\dagger N$ -type terms;
 - b) One-loop coupling constant correction (to h) due to $K^\dagger K N^\dagger N$ -type terms.
- The solid lines represent nucleons and the dotted lines represent mesons.

Figure 2: Possible flows for given values of (a, D) in the (\tilde{M}, h) plane.

a,b) $D = 0$, no coupling; a) h is *irrelevant* and b) *relevant*. No possibility for sign change of \tilde{M} .

c) $D > 0$ and $a > 0$; the interaction term is attractive and *irrelevant*. Sign change of \tilde{M} is inevitable for *some* (\tilde{M}_0, h_0) . This is the case that corresponds to the situation encountered in the chiral perturbation theory calculation of refs.[3, 4].

d) $D < 0$ and $a > 0$; the interaction term is repulsive and *irrelevant*. Sign change of $\tilde{M} > 0$ is impossible.

e) $D > 0$ and $-1 < a < 0$; the interaction term is attractive and weakly *relevant*. Sign change is inevitable for *some* (\tilde{M}_0, h_0) .

f) $D < 0$ and $a < 0$; the interaction term is repulsive and *relevant*. Sign change is impossible for *any* $\tilde{M}_0 > 0$

g) $D > 0$ and $a < -1$; the interaction term is attractive and *relevant*. Sign change is inevitable for *any* $\tilde{M}_0 > 0$

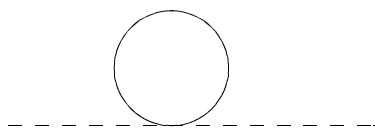
h) $D < 0$ and $a < 0$; the interaction term is repulsive and *relevant*. Sign change is impossible for *any* $\tilde{M}_0 > 0$

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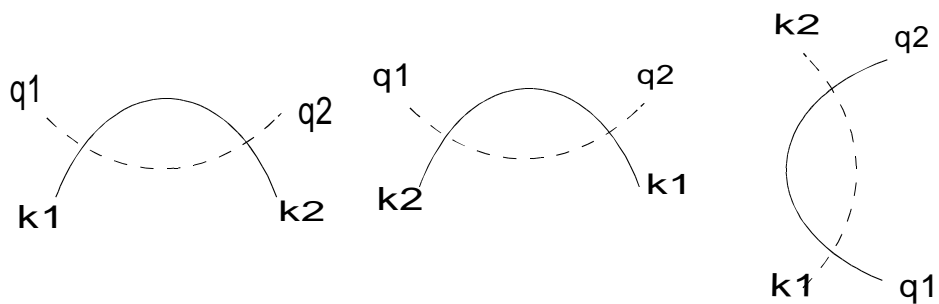
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(a)



(b)

Fig. 1

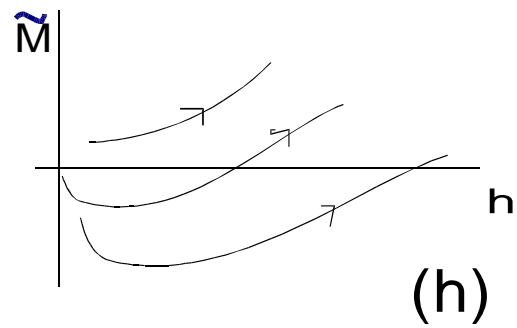
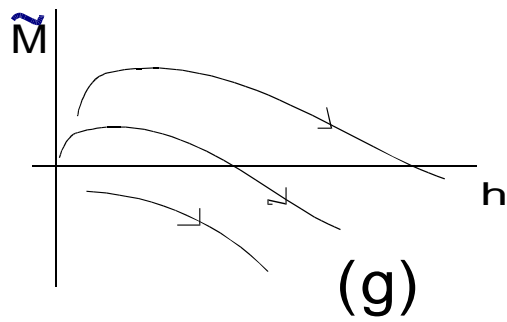
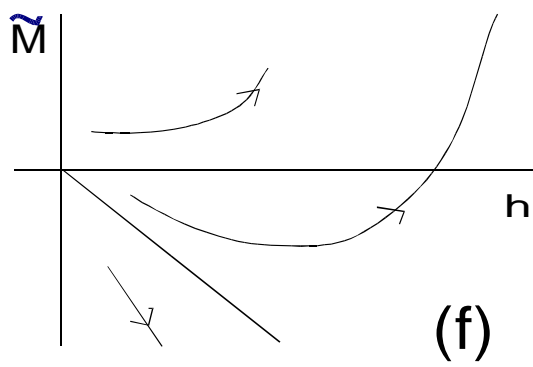
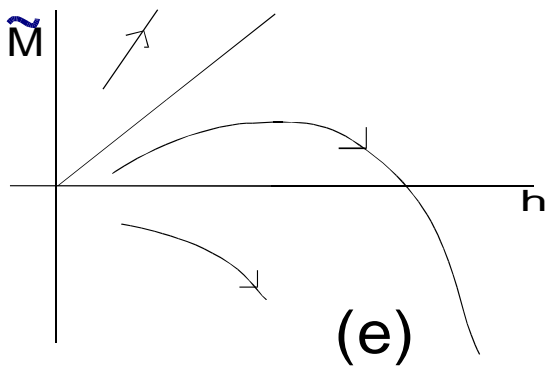
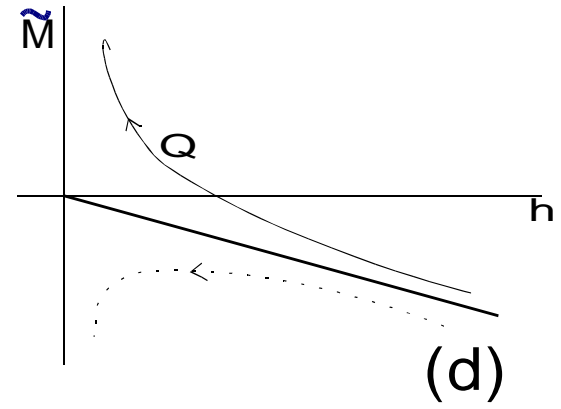
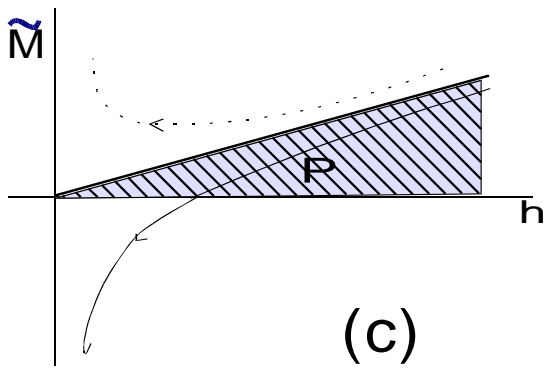
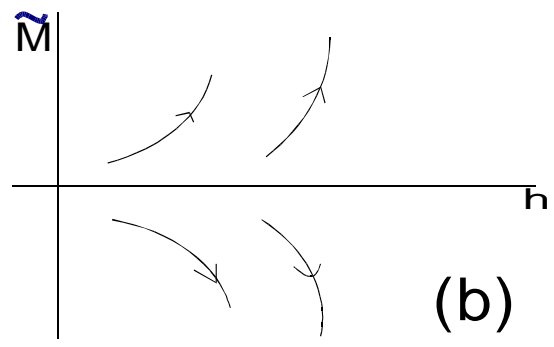
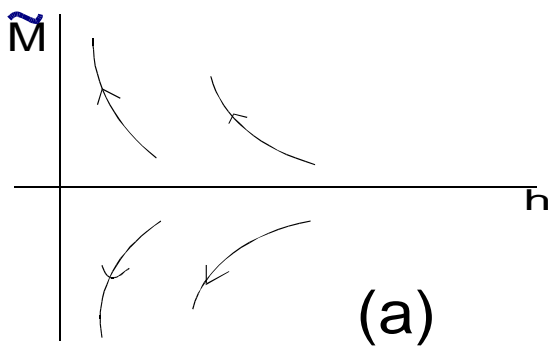


Fig. 2